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SMALL-SCALE ENERGY STORAGE AND RELEASE AS THE CAUSE OF THE STELLAR X-RAY CORONA

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Abstract. It has yet to be established why ordinary stars possess X-ray coronas. Detailed observations of the sun make it doubtful that the X-ray corona is produced by the dissipation of Alfven waves. The uniform X-ray brightness over all scales from 10^4 to 10^5 km simply does not look like a resonance phenomenon.

It is demonstrated that the arbitrary winding patterns introduced into the bipolar magnetic fields of the X-ray corona produce discontinuities within the field. It is suggested that the random motions of the footpoints of the field cause the accumulation of internal strains in the field, which dissipate through neutral point reconnection across the associated discontinuities to provide the primary heat source for the X-ray corona. It is emphasized, however, that the essential high resolution observations of the sun have yet to be carried out, and until the theory is firmly established for the sun, we cannot know how to interpret the X-ray emission of the sun or of other stars.

1. INTRODUCTION

The emission of X-rays is a universal property of stars, and a variety of ideas have been invented to explain the phenomenon. The more intense X-ray sources require matter falling on to massive compact objects to provide the basic source of energy. However, the ordinary star is more subtle. The energy source can only be the convective zone, operating as a turbulent heat engine (Alfven, 1947; Schwarzschild, 1948; Biermann, 1948) but the means by which the energy is transported to the outer atmosphere has proved elusive. It is obvious that the question can be resolved only be detailed and quantitative study. It is not sufficient to <u>detect</u> the phenomenon of stellar X-ray emission. It is necessary to <u>see</u> the phenomenon in action, to study its detailed working. The essential effects occur on scales as small as 10^2 km or less. So the sun is the essential laboratory for studying stellar X-ray emission.

It is clear from the observations of the solar corona that there are two distinct states of the stellar corona, depending upon the connectivity of the magnetic field. The open magnetic regions, in which the lines of force extend to infinity, occur where the field is relatively weak (~ 10 gauss) and the gas density is low ($\sim 10^8$ atoms/cm^3) so that there is little emitted radiation in spite of the temperatures of 1.5-2x10 °K. These "coronal holes" continually expand outward into space, producing the fast streams--and perhaps the entirety--of the solar wind, requiring an energy input of nearly 10^6 ergs/cm^2 sec (Withbroe and Noyes, 1977). The only known mechanism for supplying this energy is the emission of Alfven waves by the photospheric motions and the dissipation of the waves through phase mixing, etc. (cf. Hayvaerts and Priest, 1983; Abdelatif, 1987) over characteristic lengths of the order of 10 $R_\odot$ or more. This idea needs to be tested quantitatively of course, beginning with the measurement of the wave spectrum emitted at the photosphere.

In contrast to the open magnetic regions, the closed magnetic regions, in which the re-entrant bipolar magnetic field may be 10^2 gauss, are copious emitters of X-rays. The gas is trapped in the magnetic field and cannot escape, with the result that the temperature rises to $2\text{-}3\times10^6$ °K, and the density (from the photosphere) builds up to 10^{10} atoms/cm so that the principal energy loss is electromagnetic emission at X-ray frequencies. The energy consumption is approximately 10^7 ergs/cm^2 sec (Withbroe and Noyes 1977). Rosner, Tucker, and Vaiana (1978) have emphasized, from their analysis of the detailed observations of the structure of the closed magnetic regions - the X-ray corona - of the sun, that there is a direct connection between the strength of the magnetic field and the energy output. That is to say, as with all other aspects of stellar activity, the X-ray emission is a magnetic phenomena. The magnetic energy density in the corona is 20 or 30 times the thermal energy density.

Fortunately the observations of the sun go on to define the problem rather closely. Studies of the widths of coronal lines (Beckers and Schneeberger, 1977; Bruner, 1978; Cheng, Doschek, and Feldman, 1979) place an upper limit of about 20 km/sec on the rms fluid velocities in the line of sight in the X-ray corona, severely limiting the energy that can be transported by Alfven waves. The observations also disclose the remarkable fact (cf. Rosner, Tucker and Vaiana, 1978) that the intensity of the X-ray emission is essentially independent of the scale of the bipolar magnetic region in which it occurs, from the small X-ray bright points (10^4 km) to the normal active regions (10^5 km). These two constraints restrict the theoretical possibilities for producing (i.e. heating) the X-ray corona.

The obvious mechanism is the coronal dissipation of Alfven waves generated in the photosphere, much as in the coronal hole. The difficulty with this idea is that we expect waves with periods of the order of 10^2 sec, which have wavelengths of 2×10^5 km in the corona. Yet these waves must dissipate equally and substantially over scales of 10^4 km in the X-ray bright points and over scales of 10^5 km in the normal active regions. Some authors appeal to resonance effects (cf. Davila, 1985) and to turbulent energy cascade (Hollweg, 1984), which all requires a broad input wave spectrum, from periods of 10^2 sec down to 10 sec, in order to function equally over all scales from 10^4 to 10^5 km.

Our own view is that the scale independence of the X-ray coronal features does not look like a resonance phenomenon. Wave heating as the principal cause of the corona requires too many special circumstances. It is time to look for alternatives.

2. SMALL-SCALE STORAGE OF MAGNETIC ENERGY

Energy is continually introduced into the bipolar magnetic fields above the surface of the sun by the random wandering of the footpoints of the field (Parker, 1983b). The individual magnetic fibrils at the photosphere are kicked about by the granules at velocities of some fraction of a km/sec. The motions cannot be well ordered, with the result that any given fibril spends its time wandering among the neighboring fibrils, say with a velocity u. The flux bundle in the corona trails out behind its wandering photospheric footpoint and meanders among the neighboring flux bundles along the random path of the footpoint. The length of the path traversed by the footpoint at the photosphere after a time t is ut.

Then consider the idealized situation illustrated in Fig. 1, wherein there is a vertical uniform field e_zB_o, except for one wandering flux bundle whose upper end is anchored at a horizontal plane z=L above the surface z=0 on which the footpoint wanders at random. The region is filled with a tenuous infinitely conducting fluid. The vertical component B_o of the field in the wandering flux

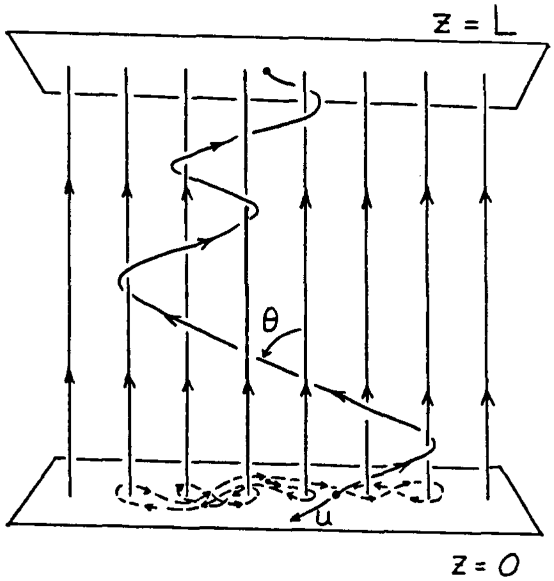

FIG. 1 A sketch of a flux bundle whose footpoint has been displaced a distance ut along a random path among the other flux bundles. The dashed line on the plane z=0 indicates the track of the footpoint.

bundle is essentially unaffected by the inclination θ of the bundle to the vertical, where

$$\tan \theta = ut/L \tag{1}$$

The horizontal component of the field is

$$\begin{aligned} B_\perp &= B_0 \tan \theta \\ &= B_0\, ut/L \end{aligned} \tag{2}$$

The tension in the inclined bundle, trailing out behind the footpoint, opposes the random walk velocity u of the footpoint with the Maxwell stress $B_\perp B_z/4\pi$, so that the motion of the footpoint does work on the field at a rate

$$\begin{aligned} W &= uB_\perp B_0/4\pi \\ &= \frac{B_0^2\, u^2 t}{4\pi L} \quad \text{ergs/cm}^2\ \text{sec} \end{aligned} \tag{3}$$

This energy accumulates in the field, of course. Every footpoint (i.e. every individual fibril) is undergoing a statistically similar random walk, with the result that W is an estimate of the average rate of energy input over the magnetic field.

Note that, beginning with a uniform vertical field at time t=0, the rate of energy input increases linearly with the passage of time, as the magnetic stresses opposing the motion of the footpoint become larger. The accumulated energy in the field is

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$$U(t) = \int_0^t dt\, W(t)$$

$$= \frac{B^2 u^2 t^2}{8\pi L} \quad \text{ergs/cm} \qquad (4)$$

and increases quadratically with time. For a normal active region with $B_0 = 10^2$ gauss and $L = 10^5$ cm, a random motion $u = 0.5$ km/sec yields

$$ut/L = 0.5\text{x}10^{-5}t, \qquad (5)$$

$$W = 2\text{x}10^2 t \text{ ergs/cm sec} \qquad (6)$$

$$U = 1\text{x}10^2 t^2 \text{ergs/cm} \qquad (7)$$

We presume that the accumulation of energy in the small-scale strains goes on until either the magnetic stress opposing the footpoint motions becomes so strong as to halt the motions, or, as seems more likely, some form of dissipation appears that destroys the accumulating small scale strains in the field as rapidly as the strains can be produced by the motion of the footpoints.

This, then, is the state of the bipolar magnetic fields in the sun. There is a substantial energy stored in small-scale strains, on scales less than 10^3 km, as a consequence of the convective motions at the photosphere.

3. SPONTANEOUS APPEARANCE OF DISCONTINUITIES

Consider how the small-scale distortions introduced in a large scale bipolar magnetic field B by the small scale ($\ell \cong 10^3$ km) motions u of the footpoints can be dissipated in a plasma at $2\text{-}3\text{x}10^6$ °K. The resistive diffusion coefficient η is only 10^3 cm /sec ($\sigma \cong 10^{17}$/sec). The characteristic dissipation time over a scale of 10^3 km is 10^5 years. There can be significant dissipation only if the characteristic gradients (curl) within the field are enormously increased above the characteristic value B/ℓ.

We are familiar with the potential fields of electrostatics and magnetostatics in nonconducting media, in which the field is described by the single scalar equation

$$\nabla^2 \phi = 0$$

There is a unique solution for any given distribution of ϕ (or $\partial\phi/\partial n$) over the boundaries, and that solution is as well behaved as the boundary conditions. The inhomogeneities introduced at the boundary decline exponentially inward from the boundary, so that gradients B/ℓ introduced at the boundaries diminish into the field.

On the other hand, a magnetic field in a highly conducting fluid of negligible pressure p satisfies the familiar force-free equation $(\nabla\text{x}\mathbf{B})\text{x}\mathbf{B}=0$, so that

$$\nabla \text{x} \mathbf{B} = \alpha(r)\mathbf{B}, \qquad (8)$$

and, with $\nabla\cdot\mathbf{B}=0$, it follows that

$$\mathbf{B}\cdot\nabla\alpha = 0\,. \qquad (9)$$

These nonlinear equations have quite different properties from the linear Laplace equation. In particular they are vector equations, so that the connectivity or topology of the frozen-in field plays an essential role. The result is tangential discontinuities in a field with any but the most special topologies. The tangential discontinuity is a surface across which $|\mathbf{B}|$ is continuous but the direction of $\mathbf{B}$ is discontinuous. The curl of the field is essentially a Dirac-delta function at the discontinuity, providing a current sheet.

The formation of current sheets in complicated field topologies has been developed from several directions. Syrovatsky (1966, 1978, 1981; Bobrova and Syrovatski, 1979; Low and Hu, 1983) pointed out that a current sheet is formed when an X-type neutral point is squashed in one direction or another. The circumstances under which the strains introduced into the field may accomplish this feat emerge from the treatment of the simple problem of a magnetic field $e_z B_0$, initially uniform, extending between the boundary planes z=0 and z=L. The footpoints at z=0 are then moved about and mixed among each other by some bounded continuous mapping with characteristic scale $\imath$ (<<L) so that the lines of force are wound and interwoven among each other in their extension from z=0 to z=L. The elementary perturbation expansion about the initial field (Parker, 1972, 1979; see also Tsinganos, 1982) produces only those solutions that are invariant along the initial field ($\partial/\partial z = 0$). Van Ballegooijen (1985) has shown recently that there is a more general perturbation solution in which a variation with z is permitted of the special form

$$\frac{\partial \alpha}{\partial z} = \frac{\partial \alpha}{\partial y}\frac{\partial A}{\partial x} - \frac{\partial \alpha}{\partial x}\frac{\partial A}{\partial y} \tag{10}$$

where α is the torsion coefficient appearing in eqns. (8) and (9), and the transverse components of the field, produced by the winding of the lines of force, are expressed in terms of the vector potential A by

$$B_x = +B_0\, \partial A/\partial y, \quad B_y = -B_0 \partial A/\partial x$$

so that

$$\alpha = -\left(\partial^2 A/\partial x^2 + \partial^2 A/\partial y^2\right)$$

Van Ballegooijen points out that equation (10) is exactly analogous to the two dimensional vorticity equation for the motion of an ideal inviscid fluid described by the stream function ψ, so that

$$v_x = +\partial\psi/\partial y, \quad v_y = -\partial\psi/\partial x$$

The vorticity is then

$$\omega = -\left(\partial^2\psi/\partial x^2 + \partial^2\psi/\partial y^2\right)$$

and the vorticity equation is

$$\frac{\partial \omega}{\partial t} = \frac{\partial \omega}{\partial y}\frac{\partial \psi}{\partial x} - \frac{\partial \omega}{\partial x}\frac{\partial \psi}{\partial y} \tag{11}$$

It follows that the magnetic field varies with z in the same way that the vorticity varies with the passage of time. There is considerable knowledge of the properties of two dimensional vortex flows of inviscid fluids (cf. Batchelor, 1967). In particular, it follows that there is a well behaved continuous solution to (10) for any bounded continuous smooth vector potential A specified on z=0. From this fact Van Ballegooijen asserted that the bounded continuous mapping of the footpoints of the field cannot produce discontinuities in the field. This overlooks the fact that the winding and interweaving of the vortex lines is restricted to the special form described by the vorticity equation (11) which, for instance, does not allow braiding of three vortex filaments), whereas we are free to introduce <u>any</u> arbitrary winding pattern through continuous motion of the footpoints of the field. It can be shown by formal development of the force-free equations (8) and (9) (Parker, 1986a,b) that discontinuities occur in α and $\mathbf{B}$ unless the mapping of the footpoints is especially restricted to the vorticity equation. The recent papers by Moffat (1985,1986) discuss the ubiquitous discontinuities (vortex sheets) as a general property of the Euler equations for the ideal inviscid fluid. Tsinganos, Distler, and Rosner (1984) develop the analogy between the lines of force and the trajectory of a Hamiltonian system in phase space.

In the present exposition we illustrate the spontaneous formation of a tangential discontinuity with a simple example (Parker, 1987). Consider a twisted flux bundle extending uniformly ($\partial/\partial z=0$) from z=0 to +L (sketched in Fig. 2) and fitting <u>smoothly</u> into the ambient uniform field $e_z B_0$ that surrounds it. That is to say, the helicity of the magnetic field declines to zero at the surface of the bundle. Indeed, we may define the surface of the bundle by the vanishing of the

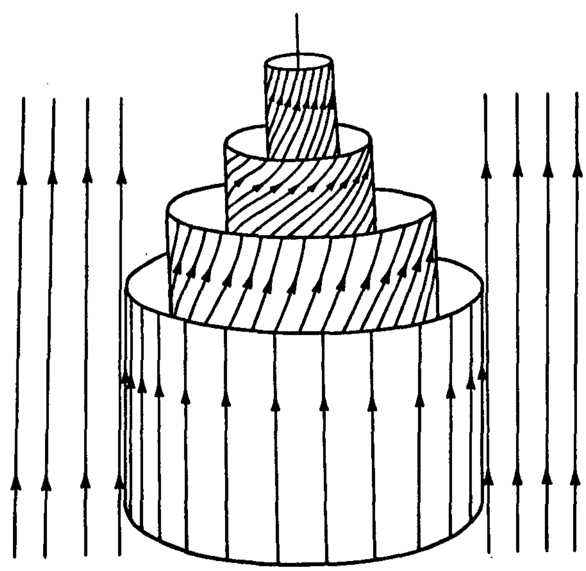

FIG. 2 A sketch of a twisted flux bundle extending uniformly from z=0 to z=L whose field fits smoothly and continuously onto the surrounding uniform field B_α.

helicity. The twisted flux bundle is in force-free equilibrium, so that the helical field within the twisted flux bundle can be expressed in terms of a generating function $f(\varpi)$, where $\varpi = (x^2 + y^2)^{\frac{1}{2}}$ represents distance from the axis of the bundle. Then (Lüst and Schlüter, 1954a,b)

$$B_z = f(\varpi) + \tfrac{1}{2}\varpi f'(\varpi), \quad B_\varphi = -\tfrac{1}{2}\varpi f'(\varpi)$$

with B falling to zero at the surface of the bundle, where $f=B_0^2$. The equilibrium field can also be described by

$$B_x = +\partial A/\partial y, \; B_y = - \partial A/\partial x, \; B_z = B_z(A)$$

with

$$\alpha = B_z'(A)$$

and
$$\frac{\partial^2 A}{\partial x^2} + \frac{\partial^2 A}{\partial y^2} + B_z(A)B_z'(A)=0$$

Formal examples of twisted flux tubes formed by bounded continuous rotation of the footpoints may be found in the literature (cf. Parker, 1987)

We consider here the case where the twisted bundle is pressed between two planes $y = \pm h$ (or two other similar twisted flux bundles) so that it is squashed into a flattened cross section, with thickness 2h and width 2w, in the manner sketched in Fig. 3. The twisted bundle resists the squashing, of course, so that some finite incremental pressure $\Delta P = \beta B_0^2/8\pi$ (where β is a number of the order of unity) must be applied by the confining planes $y = \pm h$ to squash a strongly twisted bundle. Thus, whereas the ambient pressure is $P = B_0/8\pi$, the local magnetic pressure along the midline $x=0$, $y = \pm h$ of the flattened surface of the flux tube is larger, $P + \Delta P = P(1+\beta)$. In moving in the x-direction away from the midline of the flattened side, the magnetic pressure declines by ΔP. Equilibrium requires that some force balance the outward thrust of this locally enhanced pressure.

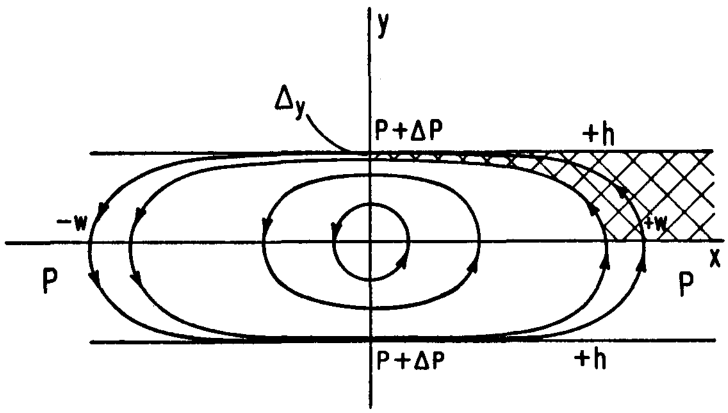

FIG. 3 A sketch of the cross section of the flux bundle in Fig. 2 after being squashed between the two planes $y = \pm h$, so that the width becomes 2w. The cross hatched region represents the thin layer of flux at the surface of the bundle, extending from the midline $x=0$ to $x= +\infty$.

Only the tension in the lines of force is available, in the amount $B_\perp^2/4\pi$ where $B_\perp$ is the transverse component of the field at the surface of the flux tube. Then the total pressure at the surface is $(B_z^2 + B_\perp^2)/8$ where B_z is the longitudinal component, from which it follows that

$$B_z^2 + B_\perp^2 = B_0^2(1+\beta) \tag{12}$$

Consider then a thin layer of flux on the surface of the twisted flux tube. The layer is indicated by the cross hatched area in Fig. 3. The thickness of the layer at the midline of the flattened side (at x=0, y = +h) is Δy, so that the tension exerted on the layer there is $\Delta y B^2/4\pi$. The enhanced pressure is ΔP, of course, so that the net force in the positive x-direction exerted on the thin layer at x=0 is

$$\Delta y(B^2/4\pi - \Delta P).$$

There is no net force in the x-direction exerted on the upper surface y=h, nor on y=0 throughout x>w, nor at x = $+\infty$ where the enhance pressure falls to zero. The only place that ΔP makes itself felt is along the lower surface of the layer. Hence the total force in the x-direction exerted on the layer is

$$F_x = \Delta y[\Delta P(0,h) - B^2(0,h)/4\pi] + \int_0^{h-\Delta y} dy\, \Delta P(x,y)$$

where the integral is along the lower surface of the layer from where it crosses the x-axis at x=w to where it crosses the y-axis at y = h-Δy. Since F_x=0 for equilibrium, it follows that

$$\frac{B^2(0,h)}{4\pi} = \Delta P(0,h) + \frac{1}{\Delta y}\int^{h-\Delta y} dy\, \Delta P(x,y) \qquad (13)$$

To estimate a lower limit on B (0,h) note that $\Delta P(x,y)>0$, so that both terms on the right hand side are positive. Hence, dropping the integral on the right hand side it follows that

$$\begin{aligned} B^2(0,h) &> 4\pi\Delta P(0,h) \\ &= \tfrac{1}{2}\beta B_o^2 \end{aligned} \qquad (14)$$

Substituting this result into (12) yields

$$B_z^2 < B_o^2\,(1+\tfrac{1}{2}\beta) \qquad (15)$$

from which it follows that

$$B_\perp/B_z > [\beta/(2+\beta)]^{\frac{1}{2}} \qquad (16)$$

The point of this inequality is that $B_\perp/B_z$ is of the order of unity, i.e. of the order of $(\beta/2)^{\frac{1}{2}}$. The field at the surface of the twisted flux bundle is inclined to the z-direction by an angle ϑ where

$$\begin{aligned} \tan\vartheta &= B\ /B_z \\ &> [\beta/(2+\beta)]^{\frac{1}{2}} \end{aligned} \qquad (17)$$

Note that the field at the surface of the flux bundle was parallel to the axis of the bundle ($B_\psi \to 0$) prior to squashing the bundle. The subsequent inclination ϑ of the surface field to the ambient field direction depends only upon the enhanced pressure $\Delta P = \beta B_o^2/8\pi$ applied to squash the tube. It is independent of the initial state of twisting of the tube. Only the sign of the inclination is obedient to the original twisting. And that is how the spontaneous tangential discontinuity is created.

Wherever a localized enhanced pressure is applied to a twisted flux bundle, the field at the surface of the bundle is strongly rotated relative to the local direction of the flux bundle . Unless the flux bundle exerting the enhanced pressure is the mirror image of the bundle being squashed, and the two bundles are precisely parallel wher they are in contact, the fields on either side of their common surface take up different directions. If, then, we shuffle and intermix the footpoints of the magnetic field in some arbitrary bounded continuous manner, the flux bundles that make up the resulting field are twisted, wound, and interwoven in arbitrary patterns that produce spontaneous misalignment of the fields across the boundaries of the local topological patterns in the winding and wrapping. The continuous fields become discontinuous because the topology becomes discontinuous where one twisted flux bundle is pressed against, or pulled around, another. The discontinuity in the topology (created when twisted tubes are pressed together) creates the discontinuity in the field direction through the rotation described by the above example[1].

This identifies, then, where the X-type neutral points may be flattened by the global magnetic stresses to form current sheets. The phenomenon arises spontaneously and unavoidably throughout any magnetic field subject to internal winding and interweaving that does not adhere exactly to the strictures of the "vorticity" equation. The interested reader is refered to the more formal treatment of the spontaneous formation of tangential discontinuities in the field (Parker, 1981a,b, 1982, 1983a,c,d, 1986a,b, 1987).

4. FORMATION OF THE STELLAR X-RAY CORONA

We suggest that the small-scale strains introduced in the bipolar magnetic fields above the surface of the sun (§II) are relieved by the dissipation that arises at the tangential discontiuities that are an essential part of the force-free equilibrium (§III). It remains for observations to establish the velocity u with which the footpoints of the field are shuffled among each other. In view of the observed horizontal motions of 1-3 km/sec at the visible surface of the sun we would imagine that u is a substantial fraction of 1 km/sec, perhaps u=0.5 km/sec.

Consider, then, the requirements for heating the X-ray corona by the dissipation at the discontinuities, presumably by neutral point reconnection in one form or another. We imagine a mean strain level in which the motion of the footpoints does work on the field at the necessary rate $W = 10^7$ ergs/cm^2 sec, and the dissipation relieves the strains as rapidly as they are introduced. Starting from an initial uniform field, it follows from eqn. (6) that W reaches 10^7 ergs/cm sec after a time $t=0.5\times10^5$ sec for the characteristic dimensions ($L=10^5$ km) of a normal active region. At this point eqn. (5) gives $ut/L = 1/4$. The local strains in the field involve inclinations to the mean field direction by a characteristic value of 14°. This same state is reached in an X-ray bright point ($L=10^4$ km) after a time $t=0.5\times10^4$ sec.

We suggest that the X-ray corona of the sun is produced primarily by this effect (Parker, 1983b). It is not possible to compute the reconnection rates associated with the power input $W = 10^7$ ergs/cm sec. We note that the rate lies between the theoretical lower and upper limits for reconnection (Parker, 1983b). It must be appreciated that the slower the reconnection rate, the more intense are the small-scale strains in the field when dissipation finally rises to the input level, and the more work is done on the field by the motions of the footpoints.

1. It is well known that a twisted flux bundle is subject to a kink instability which serves only to produce additional tangential discontinuities at the surface (Rosenbluth, Dagazian, and Rutherford, 1973).

The result is a higher level of energy input to the corona. Evidently a balance is struck when the coronal gas density rises to 10^{10} atoms/cm. and the Alfven speed (which characterizes the reconnection rate) falls to 2000 km/sec. In any case the dissipation is probably sporadic and bursty on the small-scale of the individual discontinuities (current sheets) with widths of 10^7-10^8 cm. The observed quasi-steady X-ray corona is expected to be highly active in the small.

It is obvious from this brief description of the dissipation at magnetic discontinuities that the theory of coronal heating calls upon the smallest known magnetic structures in the sun. Therefore, the theory cannot be regarded as established until observations at high spatial and spectral resolution can establish the random motions of the magnetic fibrils at the photosphère and provide direct evidence for the small-scale dissipation in the X-ray corona. Until the observations have accomplished these tasks to some satisfactory degree, the accumulating observations of stellar X-ray emission can be treated only at the phenomenological level. I hope that it will not be too many years before the theory can be established, or refuted, by observations of the sun, so that we can get on with the scientific interpretation of the fascinating stellar X-ray observations.

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